

Strong and Radiative D^* Decays

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Abstract

We use the relativistic light-front quark model to show that both strong and radiative D^* decays are in good agreement with the 1992 CLEO II results. In particular the coupling for $D^* \rightarrow D\pi$ is consistent with the experimental upper limit. The key point is the relativistic treatment of the quark spin.

1 Introduction

The D^* decays has been under both theoretical and experimental investigation for quite a long time [1, 2, 3, 4, 5, 6, 7, 8, 9]. The strong D^* decays, $D^* \rightarrow D\pi$, can be described by the lowest-order (in external momentum) effective Lagangian

$$\mathcal{L} = (\frac{m_{D^*}g}{f_\pi})\hat{D}^\dagger \boldsymbol{\tau} \cdot \partial_\mu \boldsymbol{\pi} \hat{D}^{*\mu} , \quad (1)$$

where \hat{D} and \hat{D}^* are the D and D^* isodoublet fields, and τ is the isospin matrices.

The decay width of $D^{*+} \rightarrow D^0\pi^+$, for example, reads

$$\Gamma_{D^{*+} \rightarrow D^0\pi^+} = \frac{g^2 |\mathbf{p}_\pi|^3}{6\pi f_\pi^2} . \quad (2)$$

The radiative D^* decays $D^* \rightarrow D\gamma$, on the other hand, can be described by the decay amplitude

$$M(D^* \rightarrow D\gamma) = e\mu \epsilon_{\alpha\nu\sigma\lambda} \xi^{*\alpha} \epsilon^\nu p_{D^*}^\sigma p_D^\lambda , \quad (3)$$

where $e\mu/2$ is the transition magnetic moment, and ξ^α and ϵ^ν are the polarization vectors of the photon and the D^* meson, respectively. The radiative decay width is

$$\Gamma = (e\mu)^2 \frac{|\mathbf{p}_\gamma|^3}{12\pi} . \quad (4)$$

In many theoretical studies [4, 5] people use the non-relativistic quark model to calculate the radiative D^* decays while they use various modifications of the $SU(4)$ symmetry method for the strong D^* decays. Obviously, the treatment of the strong and the radiative D^* decays is not on the same footing.

However one can also use the quark model to calculate the strong decays. Through the PCAC relation one can relate the strong coupling g to the axial-current form factor $A_0(0)$

$$g = A_0^{DD^*}(0) . \quad (5)$$

Here A_0 is one of the four form factors used to describe the matrix element of a vector meson V decaying to a pseudo-scalar meson X by the quark transition $Q' \rightarrow Q$:

$$\langle X(p_X) | \bar{Q} \gamma_\mu (1 - \gamma_5) Q' | V(p_V, \epsilon) \rangle = \frac{2V(q^2)}{m_X + m_V} i \epsilon_{\mu\nu\alpha\beta} \epsilon^{*\nu} p_X^\alpha p_V^\beta - 2m_V \frac{(\epsilon^* \cdot p_X)}{q^2} q^\mu A_0(q^2) - \left[(m_X + m_V) \epsilon^{*\mu} A_1(q^2) - \frac{(\epsilon^* \cdot p_X)}{m_X + m_V} (p_X + p_V)^\mu A_2(q^2) - 2m_V \frac{(\epsilon^* \cdot p_X)}{q^2} q^\mu A_3(q^2) \right], \quad (6)$$

where ϵ is the polarization vector of the vector meson, $q = p_V - p_X$, and

$$A_3(q^2) = \frac{m_X + m_V}{2m_V} A_1(q^2) - \frac{m_V - m_X}{2m_V} A_2(q^2), \quad A_3(0) = A_0(0). \quad (7)$$

In Eq. (5), $A_0^{DD^*}(0)$ is the $A_0(q^2 = 0)$ form factor for $D^* \rightarrow D$ induced by the quark transition $u \rightarrow d$. In the non-relativistic quark model, one obtains $g \simeq 1$ [10]. In a rather different model, the BSW model [11], one has [5]

$$g = A_0^{DD^*}(0) = \int dx \, d^2\mathbf{k}_\perp \phi_D^*(x, \mathbf{k}_\perp) \phi_{D^*}(x, \mathbf{k}_\perp), \quad (8)$$

where $\phi(x, \mathbf{k}_\perp)$ is the momentum wavefunction. Since the wavefunctions of the D and D^* are quite similar, this model also has $g \simeq 1$. In fact, in the heavy c -quark limit $\phi_{D^*}(x, \mathbf{k}_\perp)$ and $\phi_D(x, \mathbf{k}_\perp)$ should be exactly the same, due to the heavy quark symmetry, and hence in the BSW model, $g = 1$ in this limit. These results of g are, however, larger than the upper limit of g obtained from the ACCMOR measurement $\Gamma_{D^{*+}} < 131 \text{keV}$ [2]:

$$g < 0.7. \quad (9)$$

There have been some recent studies on the D^* decays [6, 7, 8, 9]. In ref. [6] the radiative D^* decays are treated with the non-relativistic quark model but the strong coupling g is related to the quark-level axial-current form factor $g_A^{ud} = \frac{3}{4}$; this is obtained from fitting the measured nucleon axial-current form factor $g_A^{\text{nucleon}} = 1.25 = \frac{5}{3} g_A^{ud}$. Thus the coupling g is $g = g_A^{ud} = \frac{3}{4}$. In refs. [7, 8] heavy-meson chiral-perturbation theory [12, 10] is used and all the D^* decays can be described by two parameters. These parameters are the couplings of the chiral Lagrangian

and can therefore be determined only by fitting the experimental data. In a recent study [9], the radiative decays are calculated using the vector–dominance hypothesis while the strong coupling g is extracted from the data on the decay $D \rightarrow \pi e \bar{\nu}_e$ using heavy–meson chiral–symmetry relations. The central value obtained for g is $g \approx 0.4$. In this study the uncertainties associated with assumptions on the f_+ form factor of $D \rightarrow \pi e \bar{\nu}_e$ and on the decay constant f_D , etc., could be large. A value of $g \sim 0.4$ is also obtained [13] in a numerical solution to another type of relativistic quark model [14].

The failure of the non–relativistic quark model and the BSW model in explaining $g < 0.7$ has its roots in the non–relativistic treatment of the quark spin in these models. In this letter, we present a calculation of the strong and radiative decays of D^* using the relativistic light–front quark model [15, 16, 17]. In our calculation we are able to treat on the same basis both the strong and radiative decays. Due to the relativistic treatment one can see unambiguously how the strong coupling g can be consistent with the limit in Eq. (9). Our calculated branching ratios for the strong and radiative decays are in good agreement with the 1992 CLEO II measurement [3].

2 The light–front relativistic quark model

The relativistic light–front quark model was developed quite a long time ago and there have been many applications [15, 16, 17]. Here we give a brief introduction.

A ground–state meson $V(\bar{Q}q)$ with spin J on the light front can be described by the state vector

$$|V(P, J_3, J)\rangle = \int d^3\mathbf{p}_1 d^3\mathbf{p}_2 \delta(\mathbf{P} - \mathbf{p}_1 - \mathbf{p}_2) \sum_{\lambda_1, \lambda_2} \Psi^{J, J_3}(\mathbf{P}, \mathbf{p}_1, \mathbf{p}_2, \lambda_1, \lambda_2) |Q(\lambda_1, \mathbf{p}_1) \bar{q}(\lambda_2, \mathbf{p}_2)\rangle, \quad (10)$$

In the light-front convention, the quark coordinates are given by

$$\begin{aligned} p_1^+ &= x_1 P^+, \quad p_2^+ = x_2 P^+, \quad x_1 + x_2 = 1, \quad 0 \leq x_{1,2} \leq 1, \\ \mathbf{p}_{1\perp} &= x_1 \mathbf{P}_\perp + \mathbf{k}_\perp, \quad \mathbf{p}_{2\perp} = x_2 \mathbf{P}_\perp - \mathbf{k}_\perp. \end{aligned} \quad (11)$$

The quantities $x_{1,2}$ and \mathbf{k}_\perp^2 are invariant under the kinematic Lorentz transformations. Rotational invariance of the wave function for states with spin J and zero orbital angular momentum requires the wave function to have the form [15, 16] (with $x = x_1$)

$$\Psi^{J,J_3}(\mathbf{P}, \mathbf{p}_1, \mathbf{p}_2, \lambda_1, \lambda_2) = R^{J,J_3}(\mathbf{k}_\perp, \lambda_1, \lambda_2) \phi(x, \mathbf{k}_\perp), \quad (12)$$

where $\phi(x, \mathbf{k}_\perp)$ is even in \mathbf{k}_\perp and

$$R^{J,J_3}(\mathbf{k}_\perp, \lambda_1, \lambda_2) = \sum_{\lambda, \lambda'} \langle \lambda_1 | R_M^\dagger(\mathbf{k}_\perp, m_Q) | \lambda \rangle \langle \lambda_2 | R_M^\dagger(-\mathbf{k}_\perp, m_{\bar{q}}) | \lambda' \rangle C^{J,J_3}(\frac{1}{2}, \lambda; \frac{1}{2}, \lambda'). \quad (13)$$

In Eq. (13), $C^{J,J_3}(\frac{1}{2}, \lambda; \frac{1}{2}, \lambda')$ is the Clebsh-Gordan coefficient and the rotation $R_M(\mathbf{k}_\perp, m_i)$ on the quark spins is the Melosh rotation [18]:

$$R_M(\mathbf{k}_\perp, m_i) = \frac{m_i + x_i M_0 - i \boldsymbol{\sigma} \cdot (\mathbf{n} \times \mathbf{k}_\perp)}{\sqrt{(m_i + x_i M_0)^2 + \mathbf{k}_\perp^2}}, \quad (14)$$

where $\mathbf{n} = (0, 0, 1)$, σ is the Pauli spin matrix, and

$$M_0^2 = \frac{m_Q^2 + \mathbf{k}_\perp^2}{x_1} + \frac{m_{\bar{q}}^2 + \mathbf{k}_\perp^2}{x_2}. \quad (15)$$

The spin wave function $R^{J,J_3}(\mathbf{k}_\perp, \lambda_1, \lambda_2)$ in Eq. (13) can also be written as

$$\begin{aligned} R^{J,J_3}(\mathbf{k}_\perp, \lambda_1, \lambda_2) &= \chi_{\lambda_1}^\dagger R_M^\dagger(\mathbf{k}_\perp, m_Q) S^{J,J_3} R_M^{\dagger T}(-\mathbf{k}_\perp, m_{\bar{q}}) \chi_{\lambda_2} \\ &= \chi_{\lambda_1}^\dagger U_V^{J,J_3} \chi_{\lambda_2}, \end{aligned} \quad (16)$$

where S^{J,J_3} is defined by

$$S^{J,J_3} = \sum_{\lambda, \lambda'} |\lambda\rangle \langle \lambda'| C^{J,J_3}(\frac{1}{2}, \lambda; \frac{1}{2}, \lambda'). \quad (17)$$

For the pseudoscalar and vector mesons,

$$S^{0,0} = \frac{i\sigma_2}{\sqrt{2}} \quad , \quad S^{1,\pm 1} = \frac{1 \pm \sigma_3}{2} \quad , \quad S^{1,0} = \frac{\sigma_1}{\sqrt{2}} \quad . \quad (18)$$

The explicit expressions of U^{1,J_3} in Eq. (16) can be found in ref. [17].

The matrix element of a vector meson $V(Q'\bar{q})$ decaying to a pseudo-scalar meson $X(Q\bar{q})$ is

$$\begin{aligned} \langle X(p_X) | \bar{Q} \Gamma Q' | V(p_V, J_3) \rangle &= \int dx \, d^2 \mathbf{k}_\perp \sum_{\lambda_1, \lambda_2} \frac{\Psi_X^{*0,0} \bar{u}_Q \Gamma u_{Q'} \Psi_V^{1,J_3}}{x} \\ &= \int dx \, d^2 \mathbf{k}_\perp \frac{\phi_X^* \phi_V}{x} \text{Tr} \left[U_X^{\dagger 0,0} U_\Gamma U_V^{1,J_3} \right] \quad , \end{aligned} \quad (19)$$

where U_Γ is defined by

$$\bar{u}_Q^i \Gamma u_{Q'}^j = \chi_i^\dagger U_\Gamma \chi_j \quad , \quad (20)$$

In Eq. (19), we choose $q^+ = 0$ so $q^2 = -\mathbf{q}_\perp^2$. In contrast to Eq. (19), the matrix element in the BSW model is given by

$$\langle X(p_X) | \bar{Q} \Gamma Q' | V(p_V, J_3) \rangle = \int dx \, d^2 \mathbf{k}_\perp \frac{\phi_X^* \phi_V}{x} \text{Tr} \left[S^{\dagger 0,0} U_\Gamma S^{1,J_3} \right] \quad . \quad (21)$$

Eq. (19) is in fact expected to be valid only for “good” currents such as $\Gamma = \gamma^+, \gamma^+ \gamma_5, \dots$. There are contributions other than the one given in Eq. (19) if a current is not a “good” current [16].

Using the “good” currents $\Gamma = \gamma^+ \gamma_5, \gamma^+$ and Eq. (19), we obtain the following expressions for the form factors $A_0(0)$ and $V(0)$ of the $Q' \rightarrow Q$ transition

$$A_0(0) = \int \frac{dx \, d^2 \mathbf{k}_\perp \phi_X^* \phi_V}{\sqrt{(A_V^2 + \mathbf{k}_\perp^2)(A_X^2 + \mathbf{k}_\perp^2)}} \left[A_V A_X + (2x - 1) \mathbf{k}_\perp^2 + \frac{2(m_{Q'} + m_Q)(1 - x) \mathbf{k}_\perp^2}{W_V} \right] \quad (22)$$

$$\begin{aligned} V(0) = \int \frac{dx \, d^2 \mathbf{k}_\perp \phi_X^* \phi_V}{\sqrt{(A_V^2 + \mathbf{k}_\perp^2)(A_X^2 + \mathbf{k}_\perp^2)}} & (m_X + m_V)(1 - x) \\ & \left[A_X + \frac{\mathbf{k}_\perp^2}{W_V} + (1 - x)(m_Q - m_{Q'}) \mathbf{k}_\perp^2 \theta_V \right] \end{aligned} \quad (23)$$

where

$$\begin{aligned} A_V &= x m_{\bar{q}} + (1-x)m_{Q'} , \quad A_X = x m_{\bar{q}} + (1-x)m_Q \\ W_V &= M_0^V + m_Q + m_{\bar{q}} , \quad \theta_V = \left(\frac{d\phi_V}{d\mathbf{k}_\perp^2}\right)/\phi_V . \end{aligned} \quad (24)$$

Here, M_0^V corresponds to Eq. (15) for the meson V .

The meson wave functions $\phi(x, \mathbf{k}_\perp)$ are model dependent and difficult to obtain; often simple forms are assumed for them. One possibility is to use the wavefunction adopted in [11]

$$\phi(x, \mathbf{k}_\perp) = N \sqrt{x(1-x)} \exp\left(-\frac{M^2}{2w^2} \left[x - \frac{1}{2} - \frac{m_Q^2 - m_{\bar{q}}^2}{2M^2}\right]^2\right) \frac{\exp\left(-\frac{\mathbf{k}_\perp^2}{2w^2}\right)}{\sqrt{\pi w^2}} , \quad (25)$$

where M is the mass of the meson. In [16] a Gaussian type of wavefunction was used,

$$\phi(x, \mathbf{k}_\perp) = N \sqrt{\frac{d\mathbf{k}_z}{dx}} \exp\left(-\frac{\mathbf{k}^2}{2\omega^2}\right) , \quad (26)$$

where \mathbf{k}_z is defined by $x_1 M_0 = E_Q + k_z$ with $E_Q = \sqrt{\mathbf{k}_\perp^2 + k_z^2 + m_Q^2}$. The parameter ω in both Eqs. (25) and (26) should be of the order of Λ_{QCD} . As we will see, the results are not too sensitive to the choice of wavefunction.

It is interesting to look at the general behavior of the wavefunction $\phi(x, \mathbf{k}_\perp)$ in the limit of $m_Q \rightarrow \infty$. The distribution amplitude $\int d^2\mathbf{k}_\perp \phi(x, \mathbf{k}_\perp)$ of a heavy meson, as is well known, should have a peak near $x \simeq x_0 = \frac{m_Q}{m_V}$. As m_Q becomes larger the width of the peak decreases and x_0 comes closer to 1. The wavefunction $\phi(x, \mathbf{k}_\perp)$ vanishes if $\mathbf{k}_\perp^2 \gg \Lambda_{\text{QCD}}^2$ and in the $m_Q \rightarrow \infty$ limit peaks at $x_0 \rightarrow 1$ as the width of $\phi(x, \mathbf{k}_\perp) \rightarrow 0$. Both wavefunctions listed have this feature.

3 The D^* decays

We now calculate the D^* decay rates in the light-front quark model. For $D^* \rightarrow D\pi$, we first look at the form factor $A_0(0)$ of the transition $Q \rightarrow Q$. In this case Eq. (22)

reduces to

$$A_0^{\bar{Q}Q}(0) = \int dx \, d^2\mathbf{k}_\perp \phi_X^* \phi_V T_{A_0} \quad (27)$$

where

$$T_{A_0} = \frac{A^2 + (2x - 1)\mathbf{k}_\perp^2 + \frac{4m_Q(1-x)\mathbf{k}_\perp^2}{W_V}}{A^2 + \mathbf{k}_\perp^2} \quad (28)$$

and

$$A = x \, m_{\bar{q}} + (1 - x) \, m_Q . \quad (29)$$

Generally $T_{A_0} < 1$ since it comes directly from the Melosh rotation. (For no Melosh rotation one would get $T_{A_0} = 1$.) How far T_{A_0} deviates from 1 depends on the relation among m_Q , $m_{\bar{q}}$ and \mathbf{k}_\perp^2 , where the average value of \mathbf{k}_\perp^2 is determined by the scale parameter in the wavefunction such as the ω in Eqs. (25) and (26). Usually as $\langle \mathbf{k}_\perp^2 \rangle$ becomes smaller T_{A_0} becomes closer to 1. If $\mathbf{k}_\perp = 0$ in Eq. (28) then $T_{A_0} \equiv 1$.

Let us consider two limiting cases for the transition quark Q . If the transition quark Q is infinitely heavy while the spectator quark is light, $x \rightarrow 1$, and therefore the third term in the numerator of T_{A_0} vanishes. Thus

$$T_{A_0} \rightarrow 1 \quad \text{and} \quad A_0^{\bar{Q}Q}(0) \rightarrow \int dx \, d^2\mathbf{k}_\perp \phi_X^* \phi_V \rightarrow 1 , \quad (30)$$

as required by the heavy quark limit. Note in this case $q^2 = q_{max}^2 = 0$. If, as in the decays $D^* \rightarrow D\pi$, the transition quark Q is light while the spectator quark is heavy, then in Eqs. (27) and (28), $x \rightarrow 0$ as $m_c \rightarrow \infty$. In this case the third term in the numerator of T_{A_0} also vanishes, but now

$$T_{A_0} \rightarrow \frac{A^2 - \mathbf{k}_\perp^2}{A^2 + \mathbf{k}_\perp^2} . \quad (31)$$

Since both \mathbf{k}_\perp^2 and A^2 are of the same order, one can expect considerable deviation from 1 for $A_0(0)$ in this case.

Obviously, the situation for $g = A_0^{DD^*}(0)$ is quite close to the second case discussed above. Thus, one can clearly see why the light-front model gives a deviation of g from 1.

To calculate the radiative decays $D^* \rightarrow D\gamma$, one can express μ in Eq. (3) in terms of the form factor $V(0)$

$$\mu = [e_c V^{\bar{c}c}(0) + e_q V^{\bar{q}q}(0)] \left(\frac{2}{m_D + m_{D^*}} \right) \quad (32)$$

where $q = u$ or d and $e_c = \frac{2}{3}$, $e_u = \frac{2}{3}$, $e_d = -\frac{1}{3}$. $V^{\bar{Q}Q}(0)$ is the $V(0)$ form factor of the transition $Q \rightarrow Q$ ($Q = c, u$ or d). $V^{\bar{Q}Q}(0)$ in the light-front model can be obtained from Eq. (23)

$$V^{\bar{Q}Q}(0) = (m_D + m_{D^*}) \int \frac{dx \, d^2\mathbf{k}_\perp \phi_D^* \phi_{D^*}}{A^2 + \mathbf{k}_\perp^2} (1-x) \left\{ A + \frac{\mathbf{k}_\perp^2}{W_{D^*}} \right\}. \quad (33)$$

To calculate $V^{\bar{c}c}(0)$, $V^{\bar{q}q}(0)$, and g , we use the two wavefunctions given in Eqs. (25) and (26) as the wavefunctions for the D and D^* mesons. We assume $m_u = m_d = m$ and expect that the light-quark masses m in the wavefunctions to be around $0.25 \sim 0.30\text{GeV}$. Similarly the scale parameter ω in both wavefunctions should be around $0.40 \sim 0.50\text{GeV}$. These values are similar to those taken in [11] and [16].

The dependence of $V^{\bar{Q}Q}(0)$ on the quark mass m_Q is similar to that in the non-relativistic quark model [19, 20]: $V^{\bar{Q}Q} = \left(\frac{m_D + m_{D^*}}{2} \right) \sqrt{\frac{m_{D^*}}{m_D}} \left(\frac{1}{m_Q} \right)$. That is $V^{\bar{Q}Q}(0)$ decreases as m_Q increases and if $m_Q \rightarrow \infty$, $V^{\bar{Q}Q}(0) \rightarrow 1$, as required by the heavy quark relation.

In the figure we show the dependence of the coupling g on the light quark mass m and ω in the wavefunctions Eqs. (25) and (26). We choose $m_c = 1.6\text{GeV}$; the coupling g has little dependence on m_c . One can see that for a fixed ω , g increases as m increases. For m between $0.25 \sim 0.30\text{GeV}$ and ω between $0.40 \sim 0.45\text{GeV}$, g is between $0.55 \sim 0.65$. To avoid adjusting parameters, we choose $m = 0.25\text{GeV}$, and $\omega = 0.4\text{GeV}$ for the calculation of decay rates. There is good agreement between the

calculated branching ratios and experiment. With this set of parameters $g = 0.6$ using Eq. (25) or Eq. (26). In the table , we show the corresponding results for $D^* \rightarrow D\pi$ and $D^* \rightarrow D\gamma$. We also show how the branching ratios change with a different choice of mass. Finally, we give the results of our calculation for the decay $D_s^{*+} \rightarrow D_s^+ \gamma$ (for $m_s = 0.40\text{GeV}$ and $m_s = 0.50\text{GeV}$). This decay is usually assumed to be the dominant mode. In a recent paper [21] it has been argued that through isospin symmetry breaking there can also be a significant correlation between the branching ratios $\text{Br}(D_s^* \rightarrow D_s \pi^0)$ and $\text{Br}(D^{*+} \rightarrow D^+ \gamma)$.

4 Conclusion

In this letter we have presented a calculation of both strong and radiative D^* decays using the relativistic light-front quark model. Our results are in good agreement with the 1992 CLEO II measurement. In particular the strong coupling g is consistent with the experimental upper limit $\Gamma_{D^{*+}} < 131\text{keV}$. The key point is the relativistic treatment of the quark spin. The fact that the relativistic treatment is essential in explaining $g < 1$ is reminiscent of the similar situation in the nucleon axial-current form factor g_A^{nucleon} where the non-relativistic quark model gives $g_A^{\text{nucleon}} = \frac{5}{3}$ and a relativistic treatment can get g_A^{nucleon} naturally down to the measured value [22].

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Figure:

The dependence of the coupling g on the light-quark mass m and ω in Eqs. (25) and (26). The solid lines are the results of Eq. (26) and the dashed lines use Eq. (25). The upper solid and dashed lines correspond to $\omega = 0.40$ GeV, and the lower solid and dashed lines are for $\omega = 0.45$ GeV.

Table. Numerical results for D^* decays. Also listed are the experimental results of CLEO II [3] and PDG [1]. All branching ratios are in %.

Decay	Br ^{a)} Eq.(25)	Br ^{a)} Eq.(26)	Br ^{b)} Eq.(25)	Br ^{b)} Eq.(26)	Br CLEO II	Br PDG
$D^{*+} \rightarrow D^0 \pi^+$	68.2	68.0	68.1	67.7	$68.1 \pm 1.0 \pm 1.3$	55 ± 4
$D^{*+} \rightarrow D^+ \pi^0$	31.6	31.5	31.6	31.4	$30.8 \pm 0.4 \pm 0.8$	27.2 ± 2.5
$D^{*+} \rightarrow D^+ \gamma$	0.3	0.5	0.4	0.9	$1.1 \pm 1.4 \pm 1.6$	18 ± 4
$D^{*0} \rightarrow D^0 \pi^0$	75.2	73.0	70.5	66.2	$63.6 \pm 2.3 \pm 3.3$	55 ± 6
$D^{*0} \rightarrow D^0 \gamma$	24.8	27.1	29.5	33.8	$36.4 \pm 2.3 \pm 3.3$	45 ± 6
	$\Gamma(\text{keV})^a)$ Eq.(25)	$\Gamma(\text{keV})^a)$ Eq.(26)	$\Gamma(\text{keV})^b)$ Eq.(25)	$\Gamma(\text{keV})^b)$ Eq.(26)		
$D^{*+} \rightarrow \text{total}$	122.8	123.0	104.9	102.1		
$D^{*0} \rightarrow \text{total}$	74.0	76.3	67.4	69.5		
$D_s^{*+} \rightarrow D_s^+ \gamma$	0.1	0.1	0.2	0.3		

a): $\omega = 0.4\text{GeV}$, $m = m_u = m_d = 0.30\text{GeV}$, $m_s = 0.50\text{GeV}$.

b): $\omega = 0.4\text{GeV}$, $m = m_u = m_d = 0.25\text{GeV}$, $m_s = 0.40\text{GeV}$.

This figure "fig1-1.png" is available in "png" format from:

<http://arXiv.org/ps/hep-ph/9406300v1>